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BOUNDS ON COMPOSITENESS FROM NEUTRINOLESS DOUBLE BETA DECAY

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Abstract

Assuming the existence of a heavy Majorana neutral particle arising from a composite model scenario we discuss the constraints imposed by present experimental limits of half-life neutrinoless double beta decay ($0\nu\beta\beta$) measurements on the coupling of the heavy composite neutrinos to the gauge bosons. For neutrino masses $M_N = 1$ TeV we obtain a rather weak lower bound on the compositeness scale: $\Lambda \geq 0.23$ TeV.

Heavy neutral Majorana particles with masses in the TeV region are predicted in various theoretical models, such as superstring-inspired E_6 grand unification [1] or left-right symmetric models [2]. In addition the possibility of a fourth generation with a heavy neutral lepton, that could be of Majorana type, is not yet ruled out [3, 4].

In this paper we discuss the possibility that a heavy Majorana neutrino might arise from a composite model of the ordinary fermions [5]. Composite models, which describe quarks and leptons as bound states of still more fundamental particles, generally called preons, have been developed as alternatives to overcome some of the theoretical problems of the standard model [6].

Although no completely consistent dynamical composite theory has been found to date, various models have been proposed, and one common, (inevitable), prediction of these models is the existence of excited states of the known quarks and leptons, much in the same way as the hydrogen atom has a series of higher energy levels above the ground state. The masses of the excited particles should not be much lower than the compositeness scale Λ, which is expected to be at least of the order of a TeV according to experimental constraints. For example the search for four-fermion contact interactions gives $\Lambda(eell) > 0.9 - 4.7$ TeV depending on the chirality of the coupling and on the lepton flavour [7, 8]. We expect therefore the heavy fermion masses to be of the order of a few hundred GeV. The CDF experiment has excluded excited quarks in the mass range $90 - 570$ GeV from $\gamma +$ jet and W + jet final states [12].

Phenomenological implications of heavy fermions have been discussed in the literature [10, 17] using weak isospin (I_W) and hypercharge (Y) conservation. Assuming that such states are grouped in $SU(2) \times U(1)$ multiplets, since light fermions have $I_W = 0, 1/2$ and electroweak gauge bosons have $I_W = 0, 1$, to lowest order in perturbation theory, only multiplets with $I_W \leq 3/2$ can be excited. Also, since none of the gauge fields carry hypercharge, a given excited multiplet can couple only to a light

multiplet with the same Y. In addition, current conservation forces the coupling of the heavy fermions to gauge bosons to be of the magnetic moment type.

We will only consider here the excited multiplet with $I_W = 1/2 \quad Y = -1$

$$\mathcal{E} = \begin{pmatrix} N \\ E \end{pmatrix} \tag{1}$$

which can couple to the light left multiplet

$$\ell_L = \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} = \frac{1-\gamma_5}{2} \begin{pmatrix} \nu \\ e \end{pmatrix} \tag{2}$$

through the gauge fields $\vec{W}^\mu$and B^μ, with the additional assumption that N is a neutral Majorana fermion.

In terms of the physical gauge fields $W^\pm_\mu = (1/\sqrt{2})\Big(W^1_\mu \mp i\, W^2_\mu\Big)$ the relevant effective interaction can be expressed as

$$\begin{aligned} \mathcal{L}_{eff} &= (\frac{gf}{\sqrt{2}\Lambda})\Big\{\Big(\overline{N}\sigma^{\mu\nu}\frac{1-\gamma_5}{2}\, e\Big)\, \partial_\nu W^+_\mu \\ &\qquad + \Big(\overline{E}\sigma^{\mu\nu}\frac{1-\gamma_5}{2}\, \nu\Big)\, \partial_\nu W^-_\mu + h.c.\Big\} + \text{neutral currents.} \end{aligned} \tag{3}$$

where f is a dimensionless coupling constant, Λ is the compositeness scale, and $\vec{\tau}$ are the Pauli $SU(2)$ matrices, and the rest of the notation is as usual in the standard model. An extension to quarks and other multiplets, with a detailed discussion of the spectroscopy of the excited particles can be found in Ref. [11].

Regarding the experimental mass limits on the heavy Majorana neutrinos from pair production, $Z \to N\bar{N}$, we have $M_N > 34.6$ GeV at 95 % c.l., which has been deduced from the Z line shape measurements [13], and which is independent of the decay modes. More stringent limits ≈ 90 GeV come from single excited neutrino production, $Z \to N\nu$, through the transition magnetic coupling, but these do depend on assumptions regarding the branching ratio of the decay channel chosen [8, 9, 13].

In practical calculations of production cross sections and decay rates of excited states, it has been customary [15, 16, 17] to assume that the dimensionless coupling f in Eq. (3) is of order unity. However if we assume that the excited neutrino

is of Majorana type, we have to verify that this choice is compatible with present experimental limits on neutrinoless double beta decay ($0\nu\beta\beta$):

$$A(Z) \to A(Z+2) + e^- + e^- \tag{4}$$

a nuclear decay, see Fig. 1, that has attracted much attention both from particle and nuclear physicists because of its potential to expose lepton number violation. More generally, it is expected to give interesting insights about certain gauge theory parameters such as leptonic charged mixing matrix, neutrino masses etc. The process in Eq.(4), which can only proceed via the exchange of a massive Majorana neutrino, has been experimentally searched for in a number of nuclear systems [18] and has also been extensively studied from the theoretical side [19, 20, 21].

We will consider here the decay

$$^{76}\mathrm{Ge} \to \, ^{76}\mathrm{Se} + 2e^- \tag{5}$$

for which we have from the Heidelberg-Moscow $\beta\beta$-experiment the recent limit [22] ($T_{1/2}$ is the half life = log2 × lifetime)

$$T_{1/2}\,(\,^{76}\mathrm{Ge} \to \, ^{76}\mathrm{Se} + 2e^-) \geq 1.95 \times 10^{24}\,\mathrm{yr} \quad 90\,\%\ \mathrm{c.l.} \tag{6}$$

In the following we estimate the constraint imposed by the above measurement on the coupling (f/Λ) of the heavy composite neutrino, as given by Eq. (3). The fact that neutrinoless double beta decay measurements might constraint composite models, was also discussed in ref.[5] but within the framework of a particular model and referring to a heavy Majorana neutrino with the usual γ_μ coupling.

The transition amplitude of $0\nu\beta\beta$ decay is calculated according to the interaction lagrangian:

$$\begin{aligned} \mathcal{L}_{int} \quad = \quad & \frac{g}{2\sqrt{2}} \left\{ \frac{f}{\Lambda} \bar{\psi}_e(x) \sigma_{\mu\nu}(1+\gamma_5)\psi_N(x)\partial^\mu W^{\nu(-)}(x) \right. \\ & \left. \qquad + \cos\theta_C J^h_\mu(x) W^{\mu(-)}(x) + h.c. \right\} \end{aligned} \tag{7}$$

where θ_C is the Cabibbo angle ($\cos\theta_C = 0.974$) and J^h_μ is the hadronic weak charged current

$$\begin{aligned} J^h_\mu(x) &= \sum_k j_\mu(k)\delta^3(\mathbf{x}-\mathbf{r}_k) \\ j_\mu(k) &= \overline{\mathcal{N}}(\mathbf{r}_k)\gamma_\mu(f_V - f_A\gamma_5)\tau_-(k)\mathcal{N}(\mathbf{r}_k) \end{aligned} \tag{8}$$

and where $\mathbf{r}_k$ is the coordinate of the k-th nucleon, $\mathcal{N} = \binom{\psi_p}{\psi_n}$ and $\tau_-(k) = (1/2)(\tau_1(k) - i\tau_2(k))$ is the step down operator for the isotopic spin, ($\vec{\tau}(k)$ is the matrice describing the isotopic spin of the k-th nucleon). We emphasize that in Eq. (7) we have a $\sigma_{\mu\nu}$ type of coupling as opposed to the γ_μ coupling so far encountered in all $0\nu\beta\beta$ decay calculations.

For simplicity, we carry out our analysis assuming that there are no additional contributions to $0\nu\beta\beta$ decay from light Majorana neutrinos, right handed currents or other heavy Majorana neutrinos originating from another source.

The transition amplitude is then

$$\begin{aligned} S_{fi} &= (cos\theta_C)^2(\frac{g}{2\sqrt{2}})^4\left(\frac{f}{\Lambda}\right)^2(\frac{1}{2})\int\frac{d^4k}{(2\pi)^4}\,d^4x\,d^4y e^{-ik\cdot(x-y)}\times \\ & \frac{1}{\sqrt{2}}(1-P_{12})\bar{u}(p_1)\sigma_{\mu\lambda}(1+\gamma_5)\frac{\not{k}+M_N}{k^2-M_N^2}(1+\gamma_5)\sigma_{\nu\rho}v(p_2)\times \\ & [F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2}e^{ip_1\cdot x}e^{ip_2\cdot y}f_A((k-p_1)^2)f_A((k+p_2)^2)\times \\ & (k-p_1)^\lambda(k+p_2)^\rho\frac{< f|J^\mu_h(x)\ J^\nu_h(y)|i >}{[(k-p_1)^2-M_W^2][(k+p_2)^2-M_W^2]} \end{aligned} \tag{9}$$

where $(1-P_{12})/\sqrt{2}$ is the antisymmetrization operator due to the production of two identical fermions, the functions $F(Z,\epsilon)$ are the well known Fermi functions [23] that describe the distorsion of the electron's plane wave due to the nuclear Coulomb field (ϵ_i are the electron's kinetic energies in units of m_ec^2),

$$F(Z,\epsilon) = \chi(Z,\epsilon)\frac{\epsilon+1}{[\epsilon(\epsilon+2)]^{1/2}} \tag{10}$$

$$\chi(Z,\epsilon) \approx \chi^{R.P.}(Z) = \frac{2\pi\alpha Z}{1-e^{-2\pi\alpha Z}} \quad \text{(Rosen-Primakoff aproximation)}$$

and the nucleon form factor,

$$f_A(q^2) = \frac{1}{(1+|\mathbf{q}|^2/m_A^2)^2} \tag{11}$$

with $m_A = 0.85$ GeV, is introduced to take into account the finite size of the nucleon, which is known to give important effects for the heavy neutrino case.

As is standard in such calculations, we make the following approximations: [19, 20]:

i) the hadronic matrix element is evaluated within the closure approximation

$$< f|J_h^\mu(x)\ J_h^\nu(y)|i > \approx e^{i(E_f - <E_n>)x_0} e^{i(<E_n> - E_i)y_0} < f|J_h^\mu(\mathbf{x})\ J_h^\nu(\mathbf{y})|i > \tag{12}$$

where $< E_n >$ is an average excitation energy of the intermediate states. This allows one to perform the integrations over k_0, x_0, y_0 in Eq. (9);

ii) neglect the external momenta p_1, p_2 in the propagators and use the long wave-length approximation $e^{-i\mathbf{p_1}\cdot\mathbf{x}} = e^{-i\mathbf{p_2}\cdot\mathbf{y}} \approx 1$;

iii) the average virtual neutrino momentum $< |\mathbf{k}| > \approx 1/R_0 = 40$ MeV is much larger than the typical low-lying excitation energies, so that, $k_0 = E_f + E_1 - < E_n >$ can be neglected relative to $\mathbf{k}$;

iv) the effect of W and N propagators can be neglected since $M_W \approx 80$ GeV is much greater than $|\mathbf{k}|$ in the region where the integrand is large, and we are interested in heavy neutrino masses $M_N \gg M_W$.

Using the same notation as in Ref. [20] we arrive at

$$\begin{aligned} S_{fi} &= (G_F \cos\theta_C)^2 \frac{f^2}{\Lambda^2} \frac{1}{2} 2\pi\delta(E_0 - E_1 - E_2) \times \\ & \quad \frac{1}{\sqrt{2}}(1 - P_{12})\bar{u}(p_1)\sigma_{\mu i}\sigma_{\nu j}(1+\gamma_5)v(p_2)[F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2} \times \\ & \quad M_N \sum_{kl} I_{ij} < f|j^\mu(k)j^\nu(l)|i > \end{aligned} \tag{13}$$

where I_{ij} is an integral over the virtual neutrino momentum, ($\mathbf{r}_{kl} = \mathbf{r}_k - \mathbf{r}_l\,, r_{kl} = |\mathbf{r}_k - \mathbf{r}_l|\,, x_{kl} = m_A r_{kl}$)

$$\begin{aligned} I_{ij} &= \frac{1}{M_N^2} \int \frac{d^3\mathbf{k}}{(2\pi)^3}\, e^{i\mathbf{k}\cdot\mathbf{r}_{kl}}\, \frac{(-k_i k_j)}{(1 + |\mathbf{k}|^2/m_A^2)^4} \\ &= \frac{1}{4\pi} \frac{m_A^4}{M_N^2} \frac{1}{r_{kl}} \left\{ -\delta_{ij} F_A(x_{kl}) + \frac{(\mathbf{r}_k)_i (\mathbf{r}_l)_j}{r_{kl}^2} F_B(x_{kl}) \right\} \end{aligned} \tag{14}$$

with:

$$
\begin{aligned}
F_A(x) &= \frac{1}{48} e^{-x} \left(x^2 + x\right) \\
F_B(x) &= \frac{1}{48} e^{-x}\, x^3
\end{aligned}
\tag{15}
$$

Since I_{ij} is a symmetric tensor, we can make the replacement $\sigma_{\mu i}\sigma_{\nu j} \rightarrow (1/2)\{\sigma_{\mu i}, \sigma_{\nu j}\} = \eta_{\mu\nu}\eta_{ij} - \eta_{i\nu}\eta_{i\mu} + i\gamma_5\epsilon_{\mu i\nu j}$. Then, using the nonrelativistic limit of the nuclear current

$$
j_\mu(k) = \begin{cases} f_V \tau_-(k) & \text{if } \mu = 0 \\ -f_A \tau_-(k)(\sigma_k)_i & \text{if } \mu = i \end{cases}
\tag{16}
$$

($\vec{\sigma}_k$ is the spin matrice of the k-th nucleon) we arrive, with straightforward algebra, at

$$
\begin{aligned}
S_{fi} &= M_{fi}\, 2\pi\, \delta(E_0 - E_1 - E_2) \\
M_{fi} &= (G_F \cos\theta_C)^2 \frac{1}{4} \frac{-1}{2\pi} \frac{f_A^2}{r_0 A^{1/3}}\, l\, < m >
\end{aligned}
\tag{17}
$$

where we have defined

$$
\begin{aligned}
l &= \frac{1}{\sqrt{2}}(1 - P_{12})\bar{u}(p_1)(1+\gamma_5)v(p_2)[F(Z+2,\epsilon_1)F(Z+2,\epsilon_2)]^{1/2} \\
< m > &= m_e \eta_N < f \,|\, \Omega \,|\, i > \\
\eta_N &= \frac{m_p}{M_N} m_A^2 \left(\frac{f}{\Lambda}\right)^2 \\
\Omega &= \frac{m_A^2}{m_p m_e} \sum_{k \neq l} \tau_-(k)\tau_-(l) \frac{R_0}{r_{kl}} \left[\left(\frac{f_V^2}{f_A^2} - \vec{\sigma}_k \cdot \vec{\sigma}_l \right) (F_B(x_{kl}) - 3F_A(x_{kl})) \right. \\
&\qquad \left. - \vec{\sigma}_k \cdot \vec{\sigma}_l\, F_A(x_{kl}) + \frac{\vec{\sigma}_k \cdot \mathbf{r}_{kl}\, \vec{\sigma}_l \cdot \mathbf{r}_{kl}}{r_{kl}^2} F_B(x_{kl}) \right]
\end{aligned}
\tag{18}
$$

and $R_0 = r_0 A^{1/3}$ is the nuclear radius ($r_0 = 1.1$ fm).

The new result here is the nuclear operator Ω which is substantially different from those so far encountered in $0\nu\beta\beta$ decays, due to the $\sigma_{\mu\nu}$ coupling of the heavy neutrino that we are considering. The decay width is obtained upon integration over the density of final states of the two-electron system

$$
d\Gamma = \sum_{final\, spins} |M_{fi}|^2\, 2\pi\delta(E_0 - E_1 - E_2) \frac{d^3\mathbf{p}_1}{(2\pi)^3\, 2E_1} \frac{d^3\mathbf{p}_2}{(2\pi)^3\, 2E_2}
\tag{19}
$$

and the total decay rate Γ can be cast in the form

$$\Gamma = (G_F \cos\theta_C)^4 \frac{(f_A)^4\, m_e^7\, |\eta_N|^2}{(2\pi)^5 r_0^2 A^{2/3}}\, f_{0\nu}(\epsilon_0, Z)\, |\Omega_{fi}|^2 \tag{20}$$

$$f_{0\nu} = \xi_{0\nu} f_{0\nu}^{R.P.} \tag{21}$$

$$f_{0\nu}^{R.P.} = |\chi^{R.P.}(Z+2)|^2 \frac{\epsilon_0}{30}(\epsilon_0^4 + 10\epsilon_0^3 + 40\epsilon_0^2 + 60\epsilon_0 + 30) \tag{22}$$

where, $\Omega_{fi} = < f|\Omega|i >$, ϵ_0 is the kinetic energy of the two electrons in units of $m_e c^2$, and $\xi_{0\nu}$ is a numerical factor that corrects for the Rosen-Primakoff approximation [20] used in deriving the analytical expression of $f_{0\nu}^{R.P.}$. For the decay considered in Eq.(5), we have [20] $\xi_{0\nu} = 1.7$ and $\epsilon_0 = 4$. The half-life is finally written as

$$T_{1/2} = \frac{K_{0\nu}\, A^{2/3}}{f_{0\nu}\, |\eta_N|^2\, |\Omega_{fi}|^2} \tag{23}$$

$$K_{0\nu} = (\log 2) \frac{(2\pi)^5}{(G_F \cos\theta_C m_e^2)^4} \frac{(m_e r_0)^2}{m_e f_A^4} = 1.24 \times\ 10^{16}\,\mathrm{yr}$$

Combining Eq. (23) with the experimental limit given for the decay considered in Eq. (5), we obtain a constraint on the dimensionless coupling f

$$|f| \leq \left(\frac{M_N \Lambda^2}{m_p m_A^2}\right)^{1/2} \left[\frac{K_{0\nu}\, A^{2/3}}{1.4 \times 10^{24}\,\mathrm{yr} \times f_{0\nu}(Z, \epsilon_0)}\right]^{1/4} \frac{1}{|\Omega_{fi}|^{1/2}} \tag{24}$$

Given the heavy neutrino mass M_N and the compositeness scale Λ, we only need to evaluate the nuclear matrix element Ω_{fi} to know the upper bound on the value of $|f|$ imposed by neutrinoless double beta decay.

The evaluation of the nuclear matrix elements was in the past regarded as the principal source of uncertainty in $0\nu\beta\beta$ decay calculations, but the recent high-statistics measurement [24] of the allowed $2\nu\beta\beta$ decay, a second order weak-interaction β decay, has shown that nuclear physics can provide a very good description of these phenomena, giving high reliability to the constraints imposed by $0\nu\beta\beta$ decay on non-standard model parameters.

Since we simply want to estimate the order of magnitude of the constraint in Eq. (24) we will evaluate the nuclear matrix element only approximatively. First of all

the expression of the nuclear operator in Eq. (18) is simplified making the following replacement

$$\frac{r^i_{kl} r^j_{kl}}{r^2_{kl}} \rightarrow < \frac{r^i_{kl} r^j_{kl}}{r^2_{kl}} > \rightarrow \frac{1}{3}\delta_{ij} \tag{25}$$

The operator Ω becomes then

$$\Omega \approx \frac{m_A^2}{m_p m_e}(m_A R_0) \sum_{k\neq l} \tau_-(k)\tau_-(l) \left(\frac{f_V^2}{f_A^2} - \frac{2}{3}\vec{\sigma}_k \cdot \vec{\sigma}_l \right) F_N(x_{kl}) \tag{26}$$

where $F_N = (1/x)(F_B - 3F_A) = (1/48)e^{-x}(x^2 - 3x - 3)$ with F_B and F_A given in Eq.(15).

Since we are interested in deriving the lowest possible upper bound on $|f|$ given by Eq. (24), let us find the maximum absolute value of the nuclear matrix element of the operator Ω in Eq.(18):

$$|\Omega_{fi}| \leq \frac{m_A^2}{m_p m_e}(m_A R_0)|F_N(\bar{x})| \left\{ \frac{f_V^2}{f_A^2}|M_F| + \frac{2}{3}|M_{GT}| \right\} \tag{27}$$

where $M_F = < f| \sum_{k\neq l} \tau_-(k)\tau_-(l)|i >$ and $M_{GT} = < f| \sum_{k\neq l} \tau_-(k)\tau_-(l)\vec{\sigma}_k \cdot \vec{\sigma}_l|i >$ are respectively the matrix elements of the Fermi and Gamow-Teller operators whose numerical values for the nuclear system under consideration are [19, 20], $M_F = 0$ and $M_{GT} = -2.56$. Inspection of the radial function F_N (for $x \geq 0$) shows that its maximum absolute value is attained at $x = 0$. In Eq. (27) we have evaluated F_N at $x = 2.28\,(r_{kl} = 0.5$ fm). This value of r_{kl} corresponds to the typical internuclear distance at which short range nuclear correlations become important [19], so that the region $x \leq 2.28$ does not give contributions to the matrix element of the nuclear operator. We thus find

$$|\Omega_{fi}| \leq 0.6 \times 10^3 \tag{28}$$

which together with Eq. (24) gives the *conservative* upper bound on $|f|$ shown in Fig. 2 as a function of M_N for $\Lambda = 1$ TeV.

In particular, we see that the choice $|f| \approx 1$ is compatible with bounds imposed by experimental limits on neutrinoless double beta decay rates. We emphasize that

our bound on $|f|$ is conservative, and an exact evaluation of the nuclear matrix element will give an even higher lower bound.

We also note that Eq. (24) can alternatively be used to give a lower bound on Λ as a function of M_N (assuming $|f| = 1$). This is shown in Fig. 3 where we can see that the lower bound on the compositeness scale coming from $0\nu\beta\beta$ decays is rather weak: $\Lambda > 0.23$ TeV at $M_N = 1$ TeV. In table I we summarize our bounds for some values of the excited Majorana neutrino mass. We remark that, as opposed to the case of bounds coming from the direct search of excited particles, our constraints on Λ and $|f|$ *do not depend* on any assumptions regarding the branching ratios for the decays of the heavy particle.

To obtain more stringent bounds, we need to improve on the measurements of $0\nu\beta\beta$ half-life. However, our bounds c.f. Eq. (24) on ($|f|$ or Λ) depend on the experimental $T_{1/2}$ lower limit only weakly ($\propto T_{1/2}^{\pm 1/4}$) so that to obtain an order of magnitude more stringent bound we need to push higher, by a factor of 10^4, the lower bound on $T_{1/2}$.

We should bear in mind, however, that the simple observation of a few $0\nu\beta\beta$ decay events, while unmistakably proving lepton number violation and the existence of Majorana neutrals, will not be enough to uncover the originating mechanism (including the one discussed here). In order to disentangle the various models, single electron spectra will be needed, which would require high statistics experiments and additional theoretical work.

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TABLE CAPTIONS

[Table I] Lower bounds on Λ with $|f| = 1$, and upper bounds on $|f|$ with $\Lambda = 1$ TeV, for different values of the heavy neutrino mass M_N.

FIGURE CAPTIONS

[Fig. 1] Schematic illustration of neutrinoless double beta decay $0\nu\beta\beta$ via the exchange of a Majorana neutrino.

[Fig. 2] Conservative upper bound on $|f|$ versus the heavy Majorana neutrino mass M_N, at $\Lambda = 1$ TeV.

[Fig. 3] Conservative lower bound on Λ versus the heavy Majorana neutrino mass M_N for $|f| = 1$.

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TABLE I

M_N(TeV)		0.6	0.8	1.0	1.2	1.4	1.6	1.8	2.0
Λ (TeV) $>$	$[\lvert f\rvert = 1]$	0.30	0.26	0.23	0.21	0.20	0.18	0.17	0.16
$\lvert f\rvert <$	$[\Lambda = 1\,\mathrm{TeV}]$	3.3	3.8	4.3	4.7	5.1	5.4	5.7	6.1

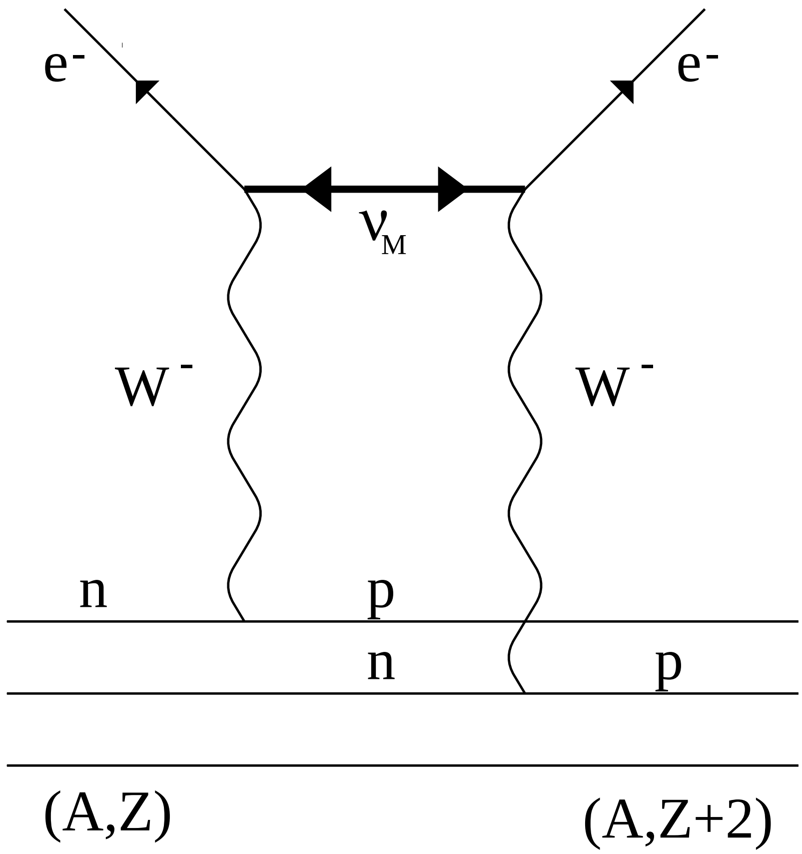

Figure 1

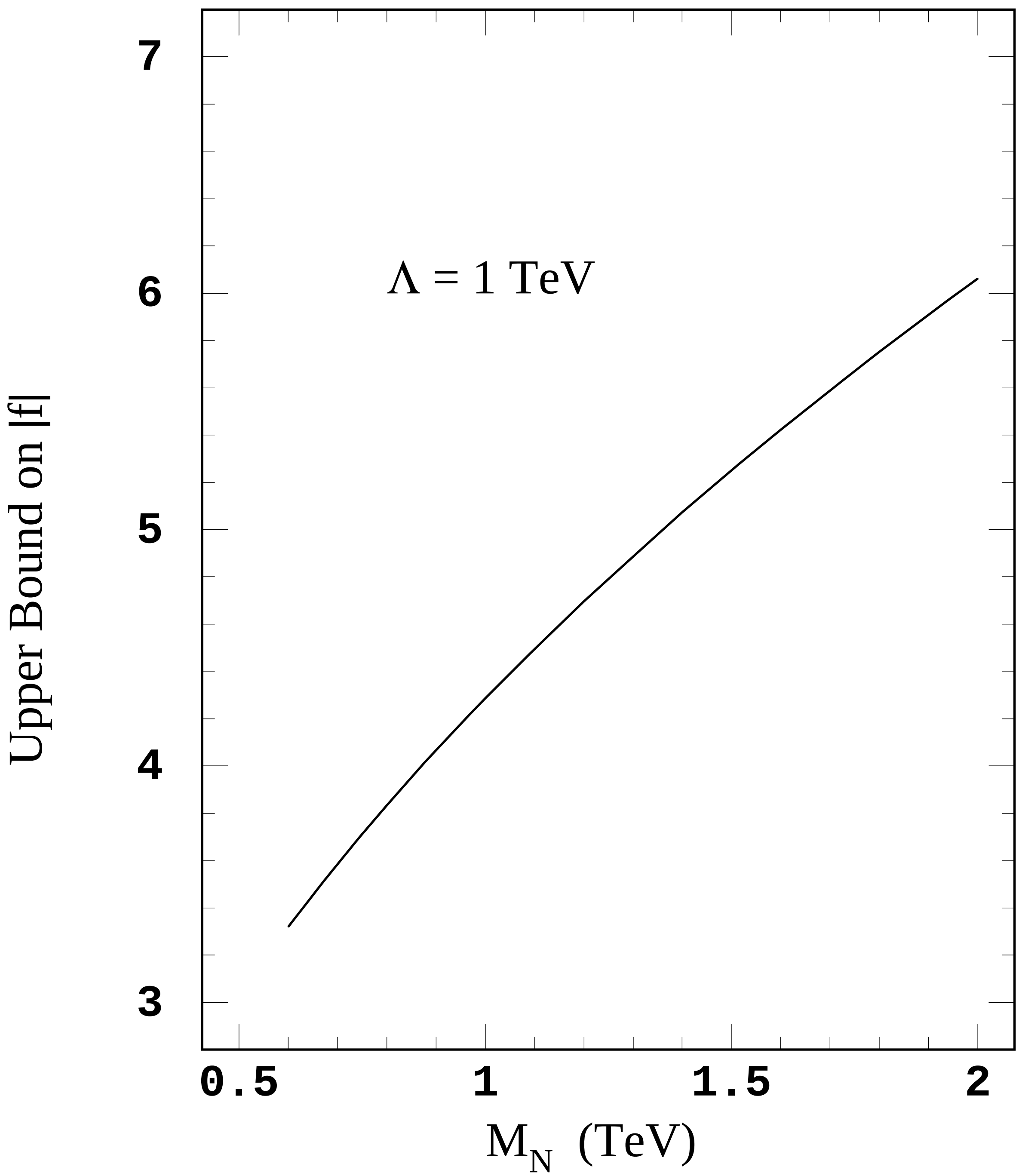

Figure 2

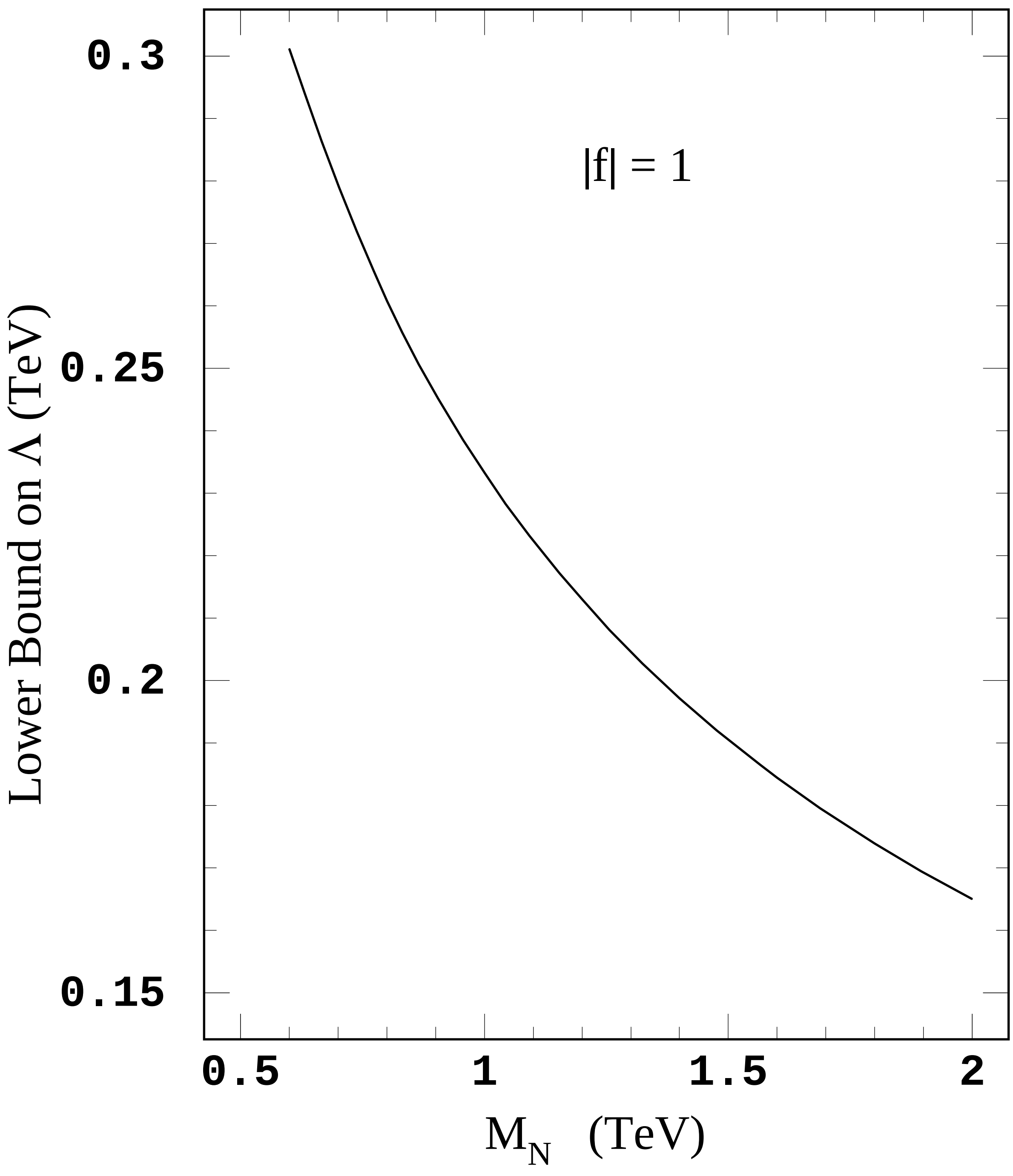

Figure 3